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DRIFT MICROWAVE PLASMA HEATING

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I. INTRODUCTION

A self-consistent plasma simulation code has been used to simulate a hot (5 keV), overdense ($\sim 10^4$), microwave-corona interaction for a hydrogen plasma.

Earlier calculations[1] have emphasized absorption during the phase of the microwave cycle when the $\tilde{B}$ field is near zero (Field Reversal Heating). These calculations, however, were not entirely self-consistent in taking into account space-charge limitations and consequently still require considerable refinement.

Another absorption mechanism occurring primarily during the electron drift portion of the cycle has been calculated self-consistently and is presented here (Drift Heating).

This code uses the electron plasma period as a time step basis, and consequently it was necessary to scale down the corona density to $\sim 10^4$ times critical from a more typical 10^8 critical (i.e., peak $\tilde{B}$ field from $>10^7$G to 10^5G). This corresponds to an incident Poynting flux of 2×10^{12} watt cm^{-2} and 10^2 plasma periods per R.f. cycle, and produces a 5 keV, $\sim 10^{17}$ cm^{-3} hydrogen plasma for which the electrons are in pressure equilibrium with the field.

The code only runs in plane coordinates, with X going into the plasma, Y tangential. The charge separation field, E_r, is along X and the incident microwave field oriented with $\tilde{B}$ along Y and $\tilde{E}$ along Z.

Since this calculation is intended to approximate an actual microwave mode field, we give a brief description of the fields appropriate to a spherical TE_{n01} mode within a spherical resonator containing a central plasma corona.[2,3] (See Figs. I & II.) A single transverse electric mode TE_{n01}, with n radial nodes, has no ϕ variation. The $\tilde{B}$ field has in general a $\hat{\theta}$ and $\hat{r}$ component, with the $\tilde{E}$ field only along $\hat{\phi}$. At the corona, $\tilde{B} = \hat{\theta} B_\theta$ only, with $B_r = 0$.

The corona radius is less than $\lambda/2$, where λ is free vacuum radiation wavelength. The B_θ field has a sin θ variation, being zero at the poles and a maximum at the waist. Then $X \to r$, $Y \to \theta$, $Z \to \phi$. Since the only field variation is with θ, on the corona radius length scale, a locally plane coordinate calculation should be an adequate first approximation.

Of more significance is the fact that the electron drift direction ϕ, along the microwave electric field, has no field variation. The only changing fields which the drifting electron sees are due solely to the intrinsic, adiabatic time dependence of the R.f. cycle (see Fig. II).

II. RESULTS

The problem was run for 3 R.f. cycles. The net absorption rate settled to approximately 2% of the incident Poynting flux, with about a 50% split between electrons and ions. The initial temperatures were 5 keV and final was $T_i = 12.4$ keV, $T_e = 13.9$ keV. (See Figs. IVa, IVb.)

An important result is that all heated particles are directed inward, towards the core.

Fig. III, a time integrated "snapshot", clearly shows the ion heating occurring primarily in bursts, twice every cycle, just after the electrons obtain their maximum energy (also twice per cycle). The electrons receive most of their energy during the phase intervals of ωt between $\pi/2 - 3\,\pi/4$ and $3\,\pi/2 - 7\,\pi/4$. That is, during the 1/8th cycle following the maximum values of $|\tilde{B}_\theta|^2$. This is also confirmed from a detailed study of the time dependent electron energy gain.

The electron and ion distribution functions are shown as Figs. IVa and IVb.

III. QUALITATIVE MODEL FOR HEATING

We modeled the microwave-corona coupling with a self-consistent orbit code which accumulates particle orbits, computes charge and current densities from them and corrects the microwave and charge separation fields to self-consistency with Maxwell's equations.

Evidently, the electrons from the core enter the R.f. skin region and drift along $\hat{\phi}$ in the crossed fields $\underset{\sim}{E}_r \times \tilde{B}_\theta$. Since the microwave field, $\hat{E}_\phi$, is along this direction, energy exchange takes place between the resonant drift electrons and the field. The electrons will continue gaining energy within the skin until the E_r field becomes small with time, such that $|E_r|/|\tilde{B}_\theta|$ no longer provides a drift concurrent with the (present) electron energy, and is reflected back into the core by the $\tilde{B}_\theta$ field.

The E_r field results from the radiation pressure of the reflected light and is quadratic in $|\tilde{B}_\theta|$. Hence E_r, like the pressure has a period half that of the R.f. cycle. The general form for E_r is known under equilibrium conditions.[4] A sketch of space and time relations between E_r, $\tilde{B}_\theta$ and $\tilde{E}_\phi$ is shown in Figs. V, VI. Note that $|\hat{E}_r| \simeq |\hat{\tilde{B}}_\theta|/2$ where $\hat{E}_r$ and $\hat{\tilde{B}}_\theta$ are peak values. This comes about because $\hat{E}_r \sim \frac{kT}{e\lambda_D}$ and for an equilibrium corona,

$$\frac{|\hat{\tilde{B}}_\theta|^2}{8\pi} \simeq 2n\,kT; \; \lambda_D = \text{Debye screening length} \simeq \left(\frac{KT}{4\pi ne^2}\right)^{\frac{1}{2}}$$

Note also that the R.f. period is adiabatic with respect to the Larmor period (in this example, $\Omega_L/2\omega \sim 10$, but for more realistic fields, $\Omega_L/2\omega \sim 10^3$). Hence an adiabatic approximation is marginal for the example, but should be good for stronger fields.

There are generally two phase regions each cycle during which drift electrons can gain energy, and two for which they can lose energy. The drift guiding center for those gaining energy tends to move such that $v_d^e/c \simeq |E_r|/|B_\theta|$ as v_d^e increases. For those losing energy, the drift will tend towards smaller values of $|E_r|/|B_\theta|$, and return quickly to the interior. (See Fig. VI.) The process then provides a net increase in electron energy, because the electrons can gain, during $\sim/10^{th}$ cycle, several times their thermal energy.

An estimate of this drift heating can then be made: During a drift time, τ, twice each cycle, electrons pick up energy ζ, with drift velocity v_d^e,

$$\zeta \simeq e\int_\tau \tilde{E}_\phi \, v_d^e(t)\, dt \simeq e\,\delta\, k\, \overline{\tilde{B}_\theta v_d^e\, \tau}$$

It might be noted that ζ is independent of $|\tilde{B}_\theta|$, for a given frequency and temperature, because $\delta \sim \omega_p^{-1} \sim n^{-\frac{1}{2}} \sim |\tilde{B}_\theta|^{-1}$, $\delta \simeq c/\omega_p$ = R.f. skin septh, $k = 2\pi/\lambda$ and where, $\tau \simeq 1/10^{th}$ cycle $= \frac{2\pi}{10\omega}$ from Fig. V.

For the example used $|\tilde{B}_\theta| \sim 10^5$ g., $\lambda \sim 1$ cm, $\delta \sim 5 \times 10^{-3}$ cm. v_θ^e at 5 keV is $\sim 3 \times 10^9$ cm. sec.$^{-1}$. Consequently, assuming a thermal velocity during drift, $\zeta \simeq 2 \times 10^{-8} \simeq 2.5$ kt. If we therefore adopt an average drift speed of $\sim 1.5\, v_\theta^e$, we obtain $\zeta \simeq 4$ kt. The average energy gain per cycle per electron is then $\frac{2 \times 4 \text{ kt}}{10} = .8$ kt. The incident flux is nv_θ^e, and rate of absorption is $\dot{E} \simeq nv_\theta^e\, .8$ kt. Since $\frac{|\tilde{B}_\theta|^2}{8\pi} \simeq 2$ n kt and the Poynting flux is $P = \frac{|\tilde{B}_\theta|^2}{8\pi} .C$, we have $\dot{E} \simeq [.4\, v_\theta^e/C]P$. This is a factor of two larger than calculated by Zohar.

When the heated drift electrons reflect back into the plasma they produce a strong additional component electrostatic field coupled to the ions, and drag ions back with them. This effect is shown in Fig. III from the code run and results in the sharing of $\dot{E}$ with the ions.

IV. MICROWAVE TARGET HEATING

The heating rate into the corona of a fuel target is $\dot{E} \simeq \text{Pf S} \frac{v_e^\theta}{c}$ where P is incident Poynting flux, S is surface area and f is an efficiency parameter, which we have seen satisfies $.2 \lesssim f \lesssim .4$ from the limited calculations made so far. A $\sim$50% split between electrons and ions is assumed. (v_θ^e is electron thermal speed.)

For a copper (or aluminum), spherical, cavity at room temperature the intrinsic quality factor $Q_o = \frac{R}{\delta_c}$, where δ_c is the R.f. skin depth, and a TE mode is assumed. At $\lambda \sim 1$ cm, $\delta_c \sim 2 \times 10^{-5}$cm.

The loaded Q is $Q_L^{-1} = \frac{1}{Q_o} + \frac{1}{Q_e}$ where Q_e represents the coupling of energy externally through the input/output guide ports. The field rise time is $\tau_c = Q_L/\omega$. For instance, with critical coupling, $Q_e = Q_o$ and $Q_L = Q_o/2$ and if R $\sim$ 10 cm, $\tau_c \simeq \frac{5 \times 10^5}{4 \times 10^{11}} \simeq 1$ µsec.

On the other hand, energy can be equilibrated between two cavities, once present, very rapidly by extreme overcoupling $Q_e \ll Q_o$, and $\tau_c \simeq$ few ns.

IV.A. RELATIVISTIC MAGNETRON GENERATED ENERGY, $\lambda \sim 1$ CM

For this system, a relativistic magnetron[6] output is coupled, near critically, into a $R \sim 10$ cm integrating cavity for $\tau_c \sim 1$ μsec, and then switched, overcoupled, to a $R \sim 10$ cm target cavity ($\tau_c^1 \sim 3$ ns). The switching is accomplished via a $\lambda/4$ wave phase shift by one of several known methods[7].

The field strength $|\tilde{B}_\theta|$ at the corona is related to that at the cavity wall, $|B_\omega|$ for the TE mode by $|B_\theta| = 2\pi \frac{R}{\lambda} |B_\omega|$.

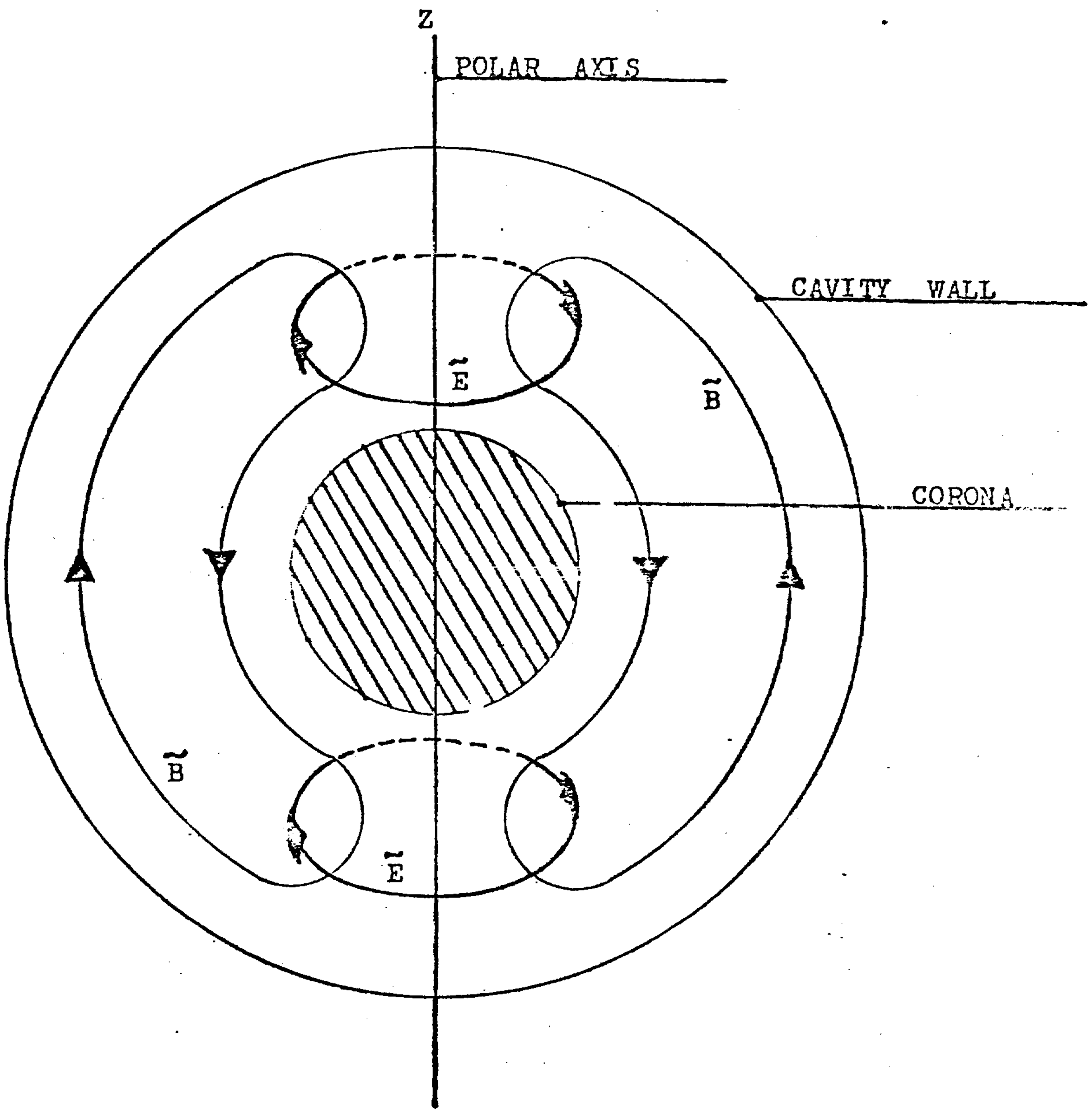
Further, the field energy, E_f, contained within the storage (and the target) cavity is $E_f = P_M \frac{Q_L}{\omega}$ where P_M is the magnetron power output. As an example, for $P_M \sim 4 \times 10^{11}$ watts, $E_f \simeq .5$ MJ. This corresponds to a $|B_\theta| \simeq 1.5 \times 10^5$ gauss which in turn provides $|B_\theta| = 10$ Mg and an incident Poynting flux $P = 10^{16}$ watts cm^{-2}. Peak wall pressure is 1 kbar $\simeq$ 15,000 psi and corona pressure is $\sim$4 Mbar.

For a corona satisfying hydroradiation stability[3], corona radius $\sim\lambda/3$, so $S \simeq 1$ cm^2. For a temperature (Corona) of 5 keV this gives $\dot{E} \simeq 2 \times 10^{14}$ for $f = .2$ or .2 MJ/ns. Since $\dot{E} \sim T_e^{1/2}$, a 20 keV temperature gives $\dot{E} \simeq 1$ MJ/ns. This system is intended for high efficiency use in power generation.

IV.B. MICROWAVE FIELD COMPRESSION

In this system, the compression cavity acts initially as the integrating cavity. By using a cryogenic temperature (initially), field levels of $\sim$1 Mg may be achievable after compression[5]. The initial frequency is at $\sim$10 cm and final frequency $\sim$3-4 mm, with energy gain of $\sim$25. For a target cavity of $R \sim 5$ cm and a $|B_\omega| \sim 1$ Mg, $|B_\theta| \simeq 70$ Mg corona pressure is $\sim$150 Mbar and an energy absorption rate $\dot{\zeta} \sim 20$ MJ/ns @ $T_e \sim 5$ keV and $f = .2$.

In either case, IV.A or IV.B, $f = .2$ corresponds, at 5 keV, to a Q of $\sim$50, and at $\lambda \sim 1$ cm this is 1.5 ns for field energy absorption. Since without the corona the cavity would absorb on a μsec time scale, virtually all the field energy in the system will go into the corona.



TE_{n01} MODE

Figure I

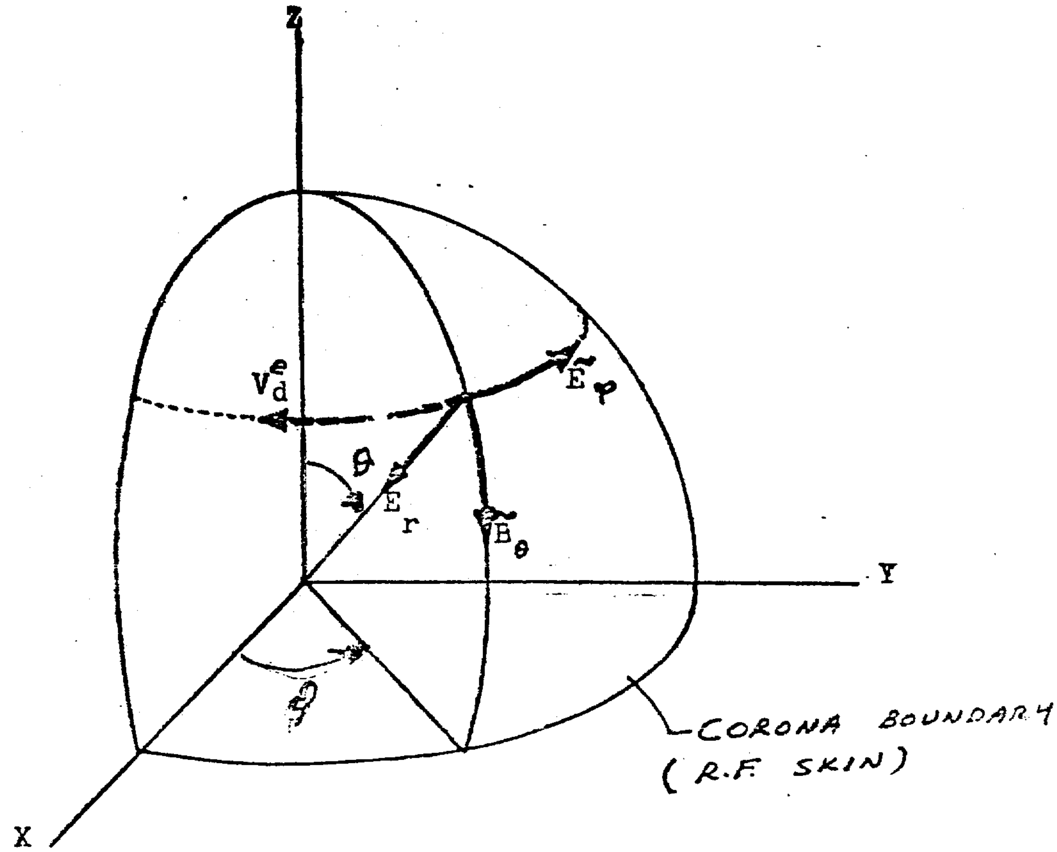

Z
V_d^e
$\tilde{E}_\varphi$
θ
$\tilde{E}_r$
$\tilde{E}_\theta$
Y
X
CORONA BOUNDARY
(R.F. SKIN)

Figure II

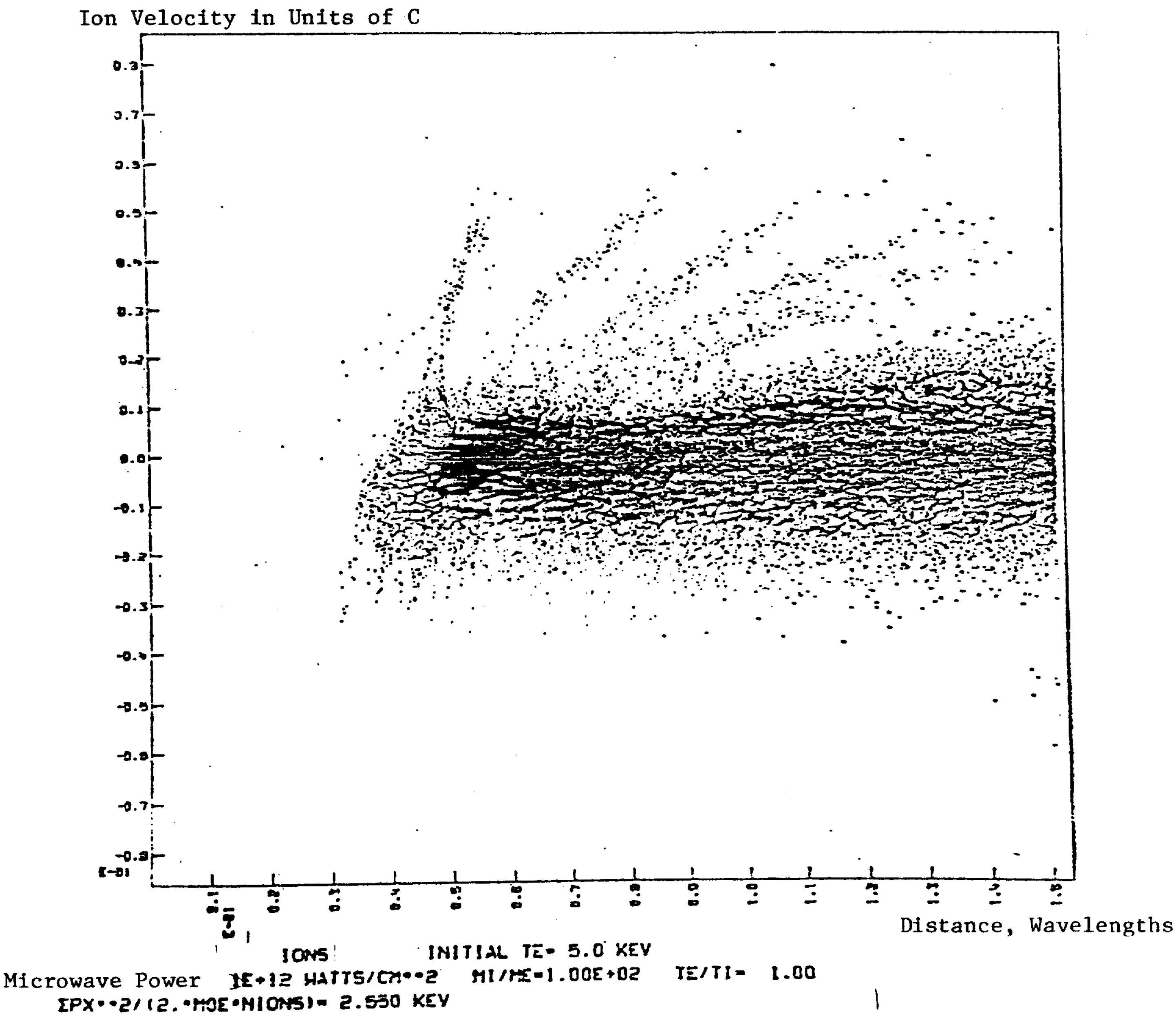
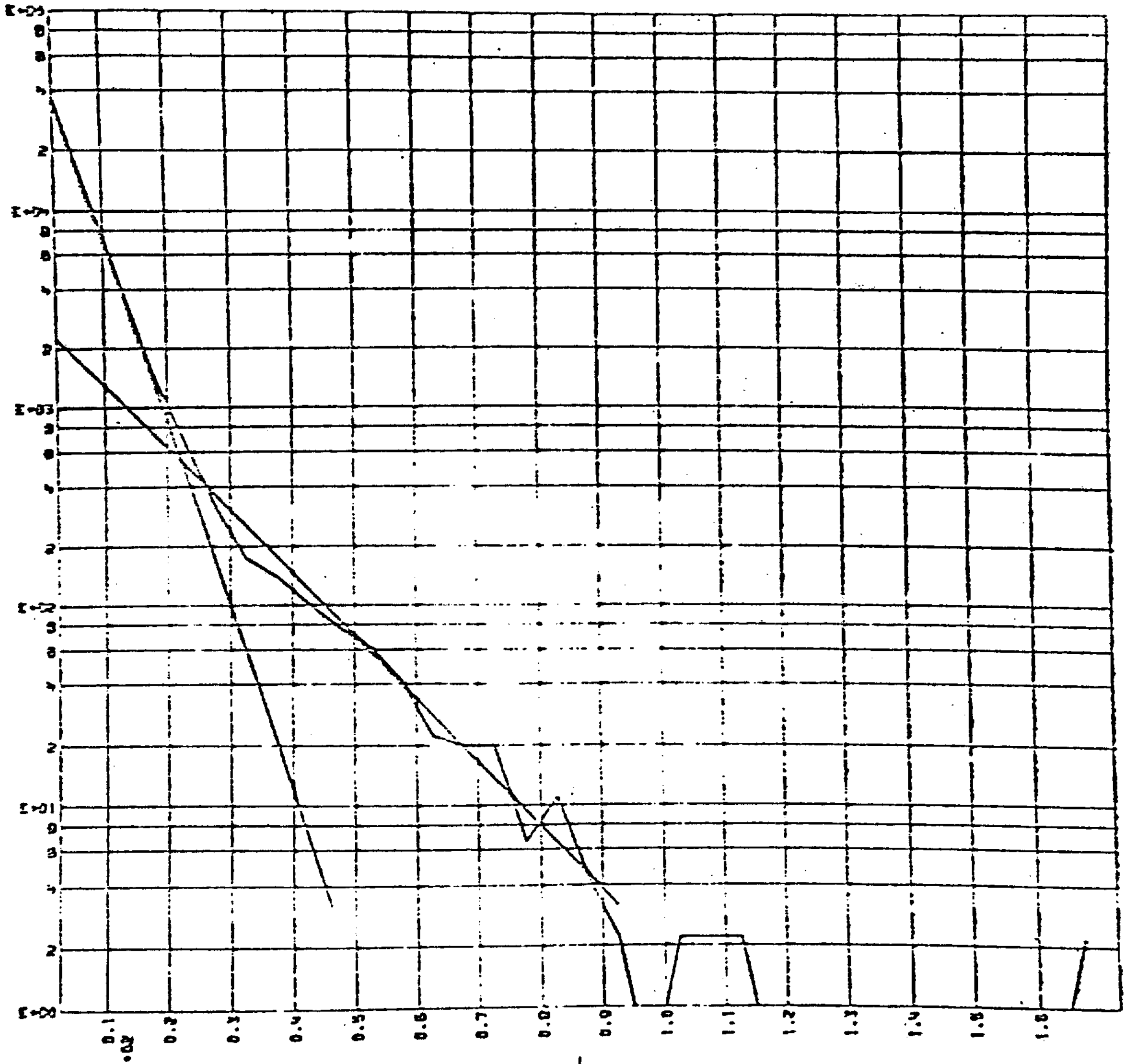


FIG. III

Ion Heating Bursts, Twice/Cycle

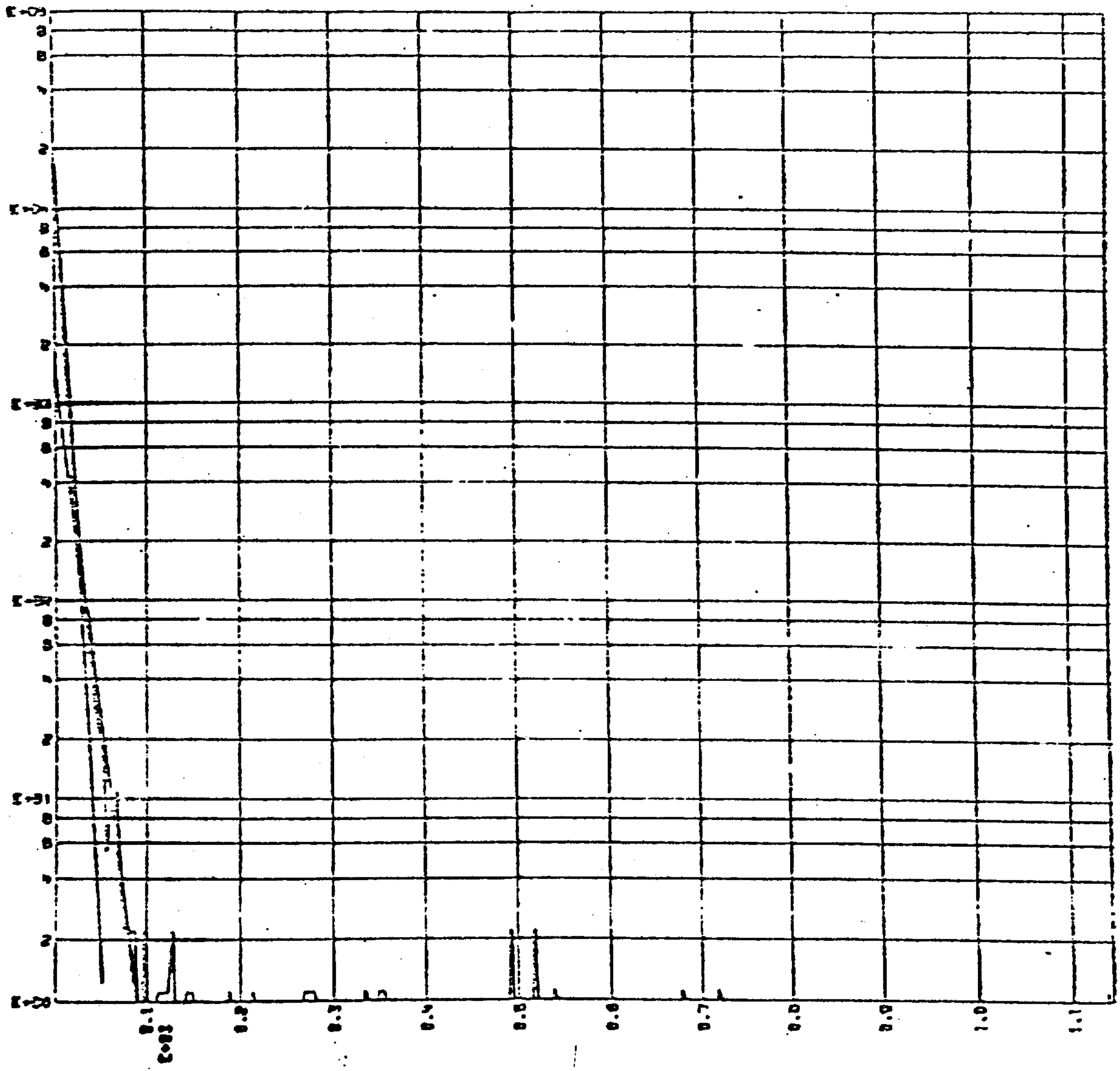


TEMPERATURES,AREAS,T*AREAS,FRACT ABS 4.69 0.84 3.95 -0.90
13.94 0.16 2.20 1.90

ION ENERGY SPECTRUM AT TIME*FO= 3.20

FIG. IVa

Ion Energy Spectrum



TEMPERATURES, AREAS, T*AREAS, FRACT ABS 5.38

12.37

ELECTRON ENERGY SPECTRUM AT TIME*F0= 3.20

FIG. IVb

Electron Energy Spectrum

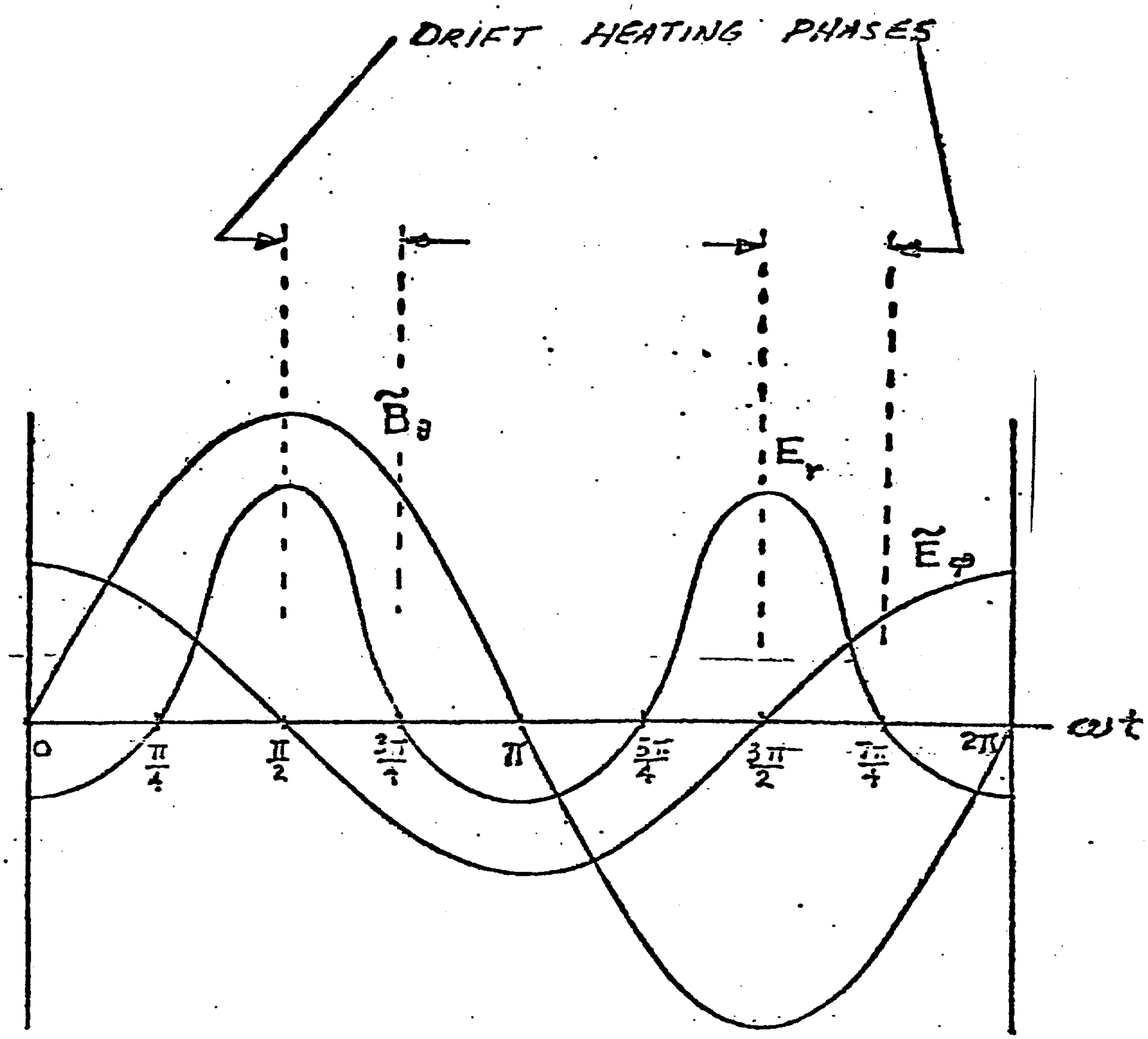

DRIFT HEATING PHASES
$\tilde{B}_\theta$
E_r
$\tilde{E}_\varphi$
ωt
0
$\frac{\pi}{4}$
$\frac{\pi}{2}$
$\frac{3\pi}{4}$
π
$\frac{5\pi}{4}$
$\frac{3\pi}{2}$
$\frac{7\pi}{4}$
2π

Figure V

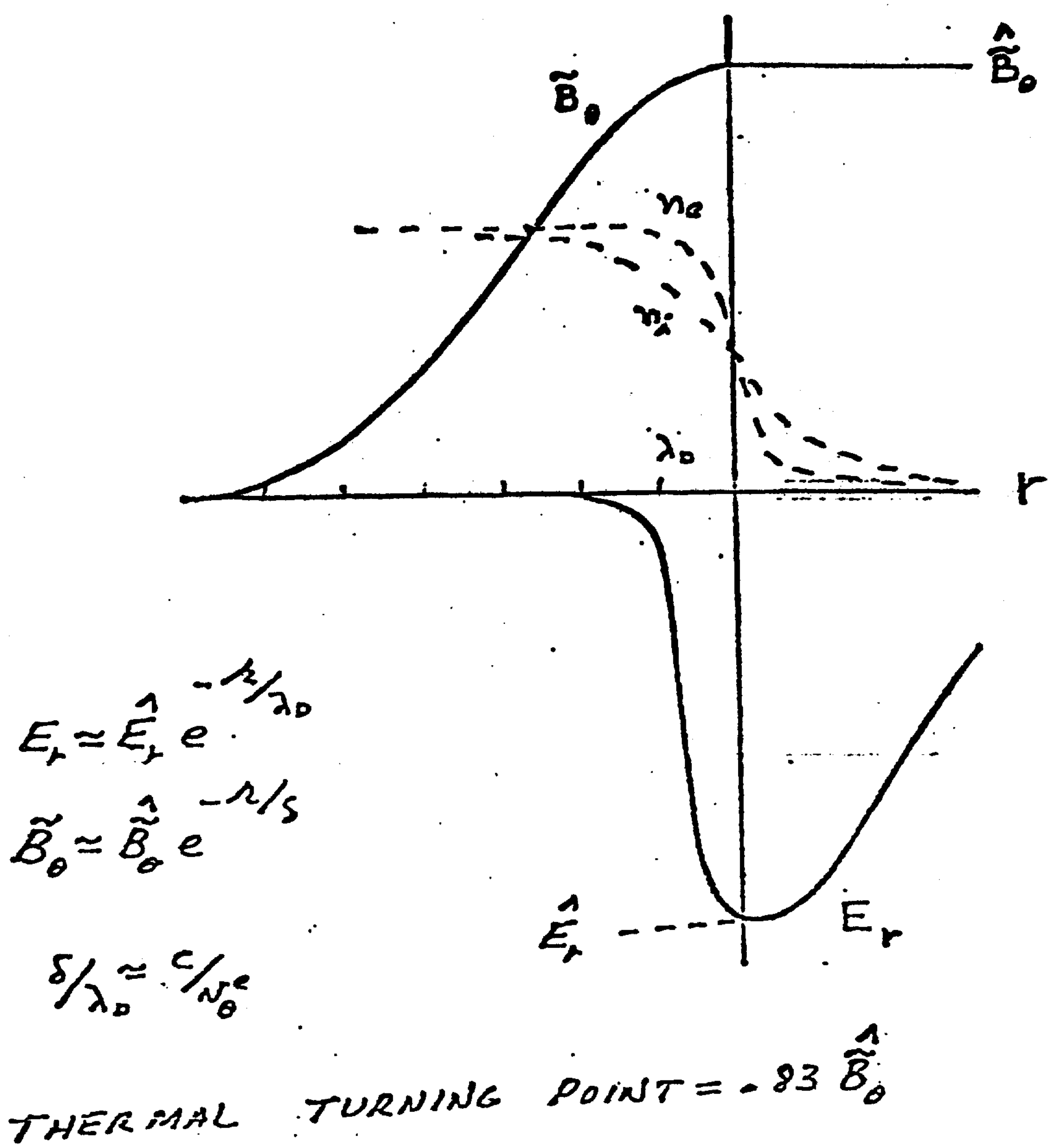


FIG. VI

Microwave Skin Region

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